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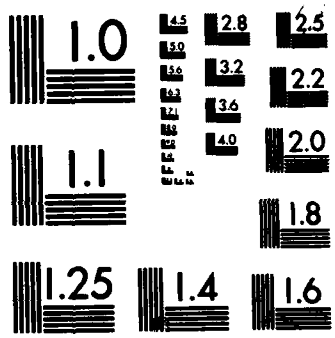

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Parametric coupling processes of a large amplitude electromagnetic pump wave with convective plasma fluid instabilities ($E \times B$ gradient drift and current convective) in the high latitude F region ionosphere have been studied. We find that a pump wave, with frequency near the upper hybrid frequency and with electric field component perpendicular to the ambient magnetic field, can stabilize or destabilize the $E \times B$ gradient-drift and/or current convective instability. For parameters typical of the nighttime high latitude F region ionosphere we find stabilization or destabilization with a vertically incident 0 mode carrying a free space incident power density on the order of 10^{-4} W/m^2.

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CONTENTS

PARAMETRIC EXCITATION AND SUPPRESSION OF CONVECTIVE PLASMA INSTABILITIES IN THE HIGH LATITUDE F-REGION IONOSPHERE

1. INTRODUCTION

Using a variety of experimental techniques, e.g., satellites [Dyson, 1969; Dyson et al., 1974; Sagalyn et al., 1974; Clark and Raitt, 1976; Phelps and Sagalyn, 1976; Rodriquez et al., 1981], rockets [Olesen et al., 1976; Ogawa et al., 1976; Kelley et al., 1980], scintillations [Aarons et al., 1973; Fremouw et al., 1977; Erukhimov et al., 1981], and radar backscatter [Weaver, 1965; Greenwald, 1974; Hower et al., 1966; Vickrey et al., 1980; Hanuise et al., 1981], it is now known that the high latitude ionosphere, from the auroral zone into the polar cap, is a highly structured and nonequilibrium medium containing irregularities (plasma density fluctuations and structures) with scale sizes ranging from hundreds of kilometers to meters. Aside from being an interesting scientific phenomenon, ionospheric irregularities are of practical interest since they can disrupt transionospheric radio wave communications channels (see recent review by Davies [1981] and references therein).

Several theories, e.g., particle precipitation, plasma instabilities and processes, and neutral fluid dynamics have been proposed to account for high latitude ionospheric irregularities (see recent review by Keskinen and Ossakow [1983a] and references therein). Recently, considerable quantitative progress has been made, especially in the area of ionospheric plasma instabilities, in identifying the physical processes that can lead to high latitude irregularities. In particular, convective plasma instabilities such as the $\underline{E} \times \underline{B}$ gradient drift instability [Simon, 1963; Linson and Workman, 1970; Keskinen and Ossakow, 1982, 1983b] and current convective instability [Lehnert, 1958; Kadomtsev and Nedospasov, 1960; Ossakow and Chaturvedi, 1979; Chaturvedi and Ossakow, 1979; Keskinen et al., 1980; Chaturvedi and Ossakow, 1981] have been invoked to explain high latitude density irregularities in and near large scale convecting auroral plasma enhancements [Vickrey et al., 1980].

However, it is well known [Silin, 1965; Dubois and Goldman, 1965; Nishikawa, 1968] that under the influence of high-frequency long wavelength electromagnetic fields electrostatic modes of oscillation in a plasma may become parametrically coupled and may grow exponentially in time or space

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before saturating at large amplitudes. These results have been used by many investigators [see reviews by Perkins et al., 1974, Fejer, 1979 and references therein] to study the role of parametric instabilities, induced by high power ground based radars, in ionospheric modification. In addition, parametric effects have been used to study the possible stabilization or destabilization of plasma instabilities, e.g., drift waves [Fainberg and Shapiro, 1967; Sundaram and Kaw, 1973] by high frequency electric fields. Lee et al. [1972] have shown that high frequency electromagnetic waves launched from ground based heaters might lead to stabilization or destabilization of the Farley-Buneman and gradient-drift instabilities in the equatorial electrojet E region ionosphere. In this paper we show that the $\underline{E} \times \underline{B}$ gradient drift and current convective instabilities in the high latitude F region ionosphere can be stabilized or destabilized by high power radio waves. In section 2 we outline the theory of the $\underline{E} \times \underline{B}$ gradient drift and current convective instabilities in the presence of a high frequency electromagnetic pump wave with frequency near the local upper hybrid frequency. We calculate the incident power density required to stabilize or destabilize these instabilities. Finally, in section 3 we summarize and discuss our results.

2. THEORY

For wavelengths greater than the ion mean free path we use fluid equations to describe the ion and electron plasma. The following geometry is used: the y-axis is in the north-south direction, the x-axis points west, and the z-axis is along the magnetic field. The set of equations, in the frame of reference in which the neutrals are at rest, become

$$\frac{\partial N_e}{\partial t} + \nabla\cdot\left(N_e \underline{V}_e\right) = 0 \tag{1}$$

$$\frac{\partial \underline{V}_e}{\partial t} + \left(\underline{V}_e\cdot\nabla\right)\underline{V}_e = -\frac{\nabla P_e}{m_e N_e} - \frac{e}{m_e}\left(\underline{E} + \frac{\underline{V}_e \times \underline{B}}{c}\right) - \nu_{en}\underline{V}_e - \underline{R}_e \tag{2}$$

$$\frac{\partial N_i}{\partial t} + \nabla\cdot\left(N_i \underline{V}_i\right) = 0 \tag{3}$$

$$\frac{\partial \underline{V}_i}{\partial t} + \left(\underline{V}_i\cdot\nabla\right)\underline{V}_i = -\frac{\nabla P_i}{m_i N_i} + \frac{e}{m_i}\left(\underline{E} + \frac{\underline{V}_i \times \underline{B}}{c}\right) - \nu_{in}\underline{V}_i - \underline{R}_i \tag{4}$$

$$\nabla\cdot\underline{E} = 4\pi e\left(N_i - N_e\right) \tag{5}$$

Here $N_\alpha(\alpha = i \text{ or } e)$ is the species density, $\underline{E} = \underline{E}_o + \underline{E}_p \cos \omega_o t - \nabla\phi$ is the total electric field which includes the ambient electric field $\underline{E}_o$, pump field $\underline{E}_p \cos \omega_o t$, and self-consistent field $\nabla\phi$. In addition, $\underline{R}_e = - \underline{R}_i = \nu_{ei}(\underline{V}_e - \underline{V}_i)$, $P_\alpha = N_\alpha T_\alpha$, ν_{ei} is the electron-ion collision frequency, $\nu_{\alpha n}$ is the electron or ion neutral collision frequency, c is the speed of light, and T_α the species temperature. Production and loss terms have been neglected. Since we will be considering perturbations with scale sizes much less than the wavelength of the pump wave we take E_p to be spatially uniform. Since pump frequencies of a few MHz corresponds to free space pump wavelengths of several hundred meters, we restrict ourselves to maximizing irregularity scale sizes on the order of tens of meters.

In equilibrium $\left(\frac{\partial}{\partial t} = 0\right)$ we have

$$\nabla \cdot N_o\left(\underline{V}_{\alpha o} + V_{\alpha p}\right) = 0 \tag{6}$$

$$0 = \pm \frac{e}{m_\alpha}\left(\underline{E}_o + \frac{\underline{V}_{\alpha o} \times B}{c}\right) - \nu_\alpha \underline{V}_{\alpha o} \tag{7}$$

$$\frac{d}{dt}\underline{V}_{\alpha p} = \pm \frac{e}{m_\alpha}\left(\underline{E}_p + \frac{\underline{V}_{\alpha p} \times B}{c}\right) - \nu_\alpha \underline{V}_{\alpha p} \tag{8}$$

where $\alpha = i,e$ and $\underline{V}_{\alpha o}$, $\underline{V}_{\alpha p}$ refer to the drifts induced by the ambient and pump electric fields, respectively. Linearizing equations (1)-(5) by assuming $N_\alpha = N_{\alpha o} + \delta n_\alpha$, $n_\alpha << N_{\alpha o}$, $\underline{V}_\alpha = \underline{V}_{\alpha o} + \delta\underline{V}_\alpha$, $\underline{E} = \underline{E}_o + \underline{E}_p \cos \omega_o t + \delta\underline{E}$ we obtain

$$\frac{\partial \delta n_\alpha}{\partial t} + \left(\underline{V}_{\alpha o} + \underline{V}_{\alpha p}\right)\cdot\nabla\delta n_\alpha + \delta\underline{V}_\alpha \cdot \nabla N_o + N_o \nabla\cdot\delta\underline{V}_\alpha = 0 \tag{9}$$

$$\frac{\partial \delta\underline{V}_\alpha}{\partial t} + \left(\underline{V}_{\alpha o} + \underline{V}_{\alpha p}\right)\cdot\nabla\delta\underline{V}_\alpha = \pm \frac{e}{m_\alpha}\left(\delta\underline{E} + \frac{\delta\underline{V}_\alpha \times B}{c}\right) - \frac{\nabla\delta p_\alpha}{m_\alpha N_o} - \nu_\alpha \delta\underline{V}_\alpha - \delta\underline{R}_\alpha \tag{10}$$

$$\nabla\cdot\delta\underline{E} = 4\pi e\left(\delta n_i - \delta n_e\right) \tag{11}$$

The equations (9)-(11) can be reduced to a single equation (dispersion relation) for one scalar function in six different ways: one may work in either the lab or oscillating frame; and one may eliminate the electron density, eliminate the ion density, or (using Poisson's equation) eliminate

both densities in favor of the electrostatic potential. We will work in the oscillating frame [Silin, 1965; Arnush et al., 1973; Lee et al., 1972] to find the density perturbations, eliminate the electron density, and finally transform back to the laboratory frame to use Poisson's equation. The oscillating frame (denoted by tildes) is defined by $\underline{\tilde{r}}_\alpha = \underline{r} - \underline{R}_\alpha(t)$, with $(d/dt)\,\underline{R}_\alpha(t) = \underline{V}_{\alpha p}(t)$, $\delta n_\alpha(\underline{r},t) = \delta\tilde{n}_\alpha(\tilde{r}_\alpha,t)$, $\delta v_\alpha(\underline{r},t) = \delta\tilde{v}_\alpha(\tilde{r}_\alpha,t)$, $\delta\underline{E}(\underline{r},t) = \delta\tilde{E}(\underline{\tilde{r}}_\alpha,t)$. Equations (9) and (10) in the oscillating frame are identical to those in the lab frame except for the substitutions $\underline{V}_{\alpha o} + \underline{V}_{\alpha p} \to \underline{V}_{\alpha o}$, $\delta n_\alpha \to \delta\tilde{n}_\alpha$, $\delta\underline{v}_\alpha \to \delta\tilde{v}_\alpha$, $\delta\underline{E} \to \delta\underline{\tilde{E}}$, $\underline{r} \to \underline{\tilde{r}}_\alpha$. Assuming $\delta\tilde{v}_\alpha$, $\delta\tilde{n}_\alpha$, $\delta\underline{\tilde{E}} \propto \exp[i(\underline{k}\cdot\underline{r} - \omega t)]$ with $\underline{k} = k_x\,\underline{\hat{x}} + k_y\,\underline{\hat{y}} + k_z\,\underline{\hat{z}}$, $\omega = \omega_r + i\gamma$, $kL \gg 1$, $L^{-1} \equiv N_o^{-1}(\nabla N_o)$, one can then define susceptibilities χ_i and χ_e by $\tilde{\delta n}_i = (-ik\chi_i/e)\delta\tilde{E}$ and $\tilde{\delta n}_e = (ik\chi_e/e)\tilde{\delta E}$ giving

$$\chi_i = -\frac{i}{4\pi}\frac{\omega_{pi}^2}{\omega-\bar{\omega}_i}\frac{1}{\Omega_i}\left[\frac{\nu_{in}}{\Omega_i}\frac{k_\perp^2}{k^2} + \frac{\Omega_i}{\nu_{in}}\frac{k_z^2}{k^2} + i\,\frac{\tilde{z}\times\underline{k}_\perp\cdot\nabla N_o}{N_o k^2}\right] \tag{12}$$

and

$$\chi_e = \frac{i}{4\pi}\frac{\omega_{pe}^2}{\omega-\bar{\omega}_e}\frac{1}{\Omega_e}\left[\frac{\tilde{\kappa}_e}{1+\tilde{\kappa}_e^2}\frac{k_\perp^2}{k^2} - \frac{\Omega_e}{\tilde{\nu}_e}\frac{k_z^2}{k^2} + i\,\frac{\tilde{\kappa}_e^2}{1+\tilde{\kappa}_e^2}\,\frac{\tilde{z}\times\underline{k}_\perp\cdot\nabla N_o}{N_o k^2}\right] \tag{13}$$

where $\omega_{p\alpha}^2 = 4\pi N_o e^2/m_\alpha$, $\bar{\omega}_i = \frac{c}{B}\,\underline{E}_{o\perp} \times \underline{\hat{z}}\cdot\underline{k} + \underline{k}\cdot\underline{V}_{oiz} + i(D_\perp k_\perp^2 + D_\parallel k_\parallel^2)$, $D_\perp = (\nu_{in}/\Omega_i)(T_i/m_i\Omega_i) + (\nu_{ei}/\Omega_e)(C_s^2/\Omega_i)$, $D_\parallel = C_s^2/\nu_{in}$, $\Omega_\alpha = eB/m_\alpha c$, $\bar{\omega}_e = \frac{c}{B}\,\underline{E}_{o\perp} \times \underline{\hat{z}}\cdot\underline{k} + \underline{k}\cdot\underline{V}_{oez}$, $\tilde{\kappa}_e = \Omega_e/\tilde{\nu}_{ei}$, $\tilde{\nu}_{ei} = \nu_{ei} - i\omega$, and $k_\perp^2 = k_x^2 + k_y^2$. In deriving the above susceptibilities we have used $\Omega_e/\nu_{ei} \gtrsim \Omega_i/\nu_i \gg 1$ (F-region approximation). In addition, we have modeled weak cold diffuse auroral field aligned currents as a relative drift, between electrons and ions, with velocity $\underline{V}_d = \underline{\hat{z}}V_d = \tilde{z}(V_{oez} - V_{oiz})$.

Having found the susceptibilities in the oscillating frame we must now transform back to the laboratory frame to use Poisson's eqn. (5). (Since we will be considering high frequency pump waves with $\omega_o \gg \omega_{pi}$ one can neglect the effect of the pump on the ions.) This transformation and resultant dispersion relation can be found in Arnush et al. [1973] and Lee et al. [1972] and can be written:

$$[1 + 4\pi(\chi_i + \chi_e)] = -J_1^2(\zeta)4\pi\,\chi_i(1 + 4\pi\,\chi_e)\;\times$$

$$\left[\frac{1}{H_e(\omega+\omega_o)} + \frac{1}{H_e(\omega-\omega_o)}\right] \tag{14}$$

where $J_n(\zeta) = \frac{i^{-n}}{\pi} \int_0^{\pi} e^{i\zeta\cos\theta} \cos n\theta \, d\theta$ is the Bessel function of integral order n, ζ is defined by $\underline{k} \cdot \underline{R} = \zeta \sin(\omega_o t + \beta)$, and $H_e = -(4\pi\chi_e)^{-1}(1 + 4\pi\chi_e)$. In deriving (14), the approximation $\zeta \ll 1$ has been made, i.e., the electron excursion length is small compared to the perturbation wavelength.

In the absence of the pump $E_p = 0$, for low frequency modes, $\delta n_e \approx \delta n_i$. This gives

$$(\omega - \bar{\omega}_e)\left[\frac{\nu_{in}}{\Omega_i}\frac{k_\perp^2}{k^2} + \frac{\Omega_i}{\nu_{in}}\frac{k_z^2}{k^2} + i\frac{\hat{z} \times \underline{k}_\perp \cdot \nabla N_o}{N_o k^2}\right] = (\omega - \bar{\omega}_i)\left[\frac{\nu_{ei}}{\Omega_e}\frac{k_\perp^2}{k^2} - \frac{\Omega_e}{\nu_{ei}}\frac{k_z^2}{k^2} + i\frac{\hat{z} \times \underline{k}_\perp \cdot \nabla N_o}{N_o k^2}\right] \tag{15}$$

with solution $\omega = \omega_{\underline{k}r}^L + i\gamma_k^L$ given by the linear growth rate [Keskinen and Ossakow, 1982]

$$\omega_{\underline{k}r}^L = \frac{c}{B} \underline{E}_{o\perp} \times \hat{\underline{z}} \cdot \underline{k} \tag{16a}$$

$$\gamma_{\underline{k}}^L = \left(\frac{k_z^2}{k_\perp^2} + \frac{\nu_{in}}{\Omega_i}\frac{\nu_{ei}}{\Omega_e}\right)^{-1} \frac{\nu_{ei}}{\Omega_i}\left(\frac{\nu_{in}}{\Omega_i}\frac{c}{B}\underline{k}_\perp \cdot \underline{E}_o - k_z V_d\right)\frac{\hat{\underline{z}} \times \underline{k}_\perp \cdot \nabla N_o}{k_\perp^2 N_o} - D_\perp k_\perp^2 - D_\parallel k_z^2 \tag{16b}$$

with $D_\perp \approx \frac{\nu_{in}}{\Omega_i}\frac{T}{m_i \Omega_i} + \frac{\nu_{ei}}{\Omega_e}\frac{C_s^2}{\Omega_i}$ and $D_\parallel \approx \frac{C_s^2}{\nu_{in}}$.

For finite E_p we have

$$(\omega - \bar{\omega}_e)\left[\frac{\nu_{in}}{\Omega_i}\frac{k_\perp^2}{k^2} + \frac{\Omega_i}{\nu_{in}}\frac{k_z^2}{k^2} + i\frac{\hat{z} \times \underline{k}_\perp \cdot \nabla N_o}{N_o k^2}\right] - (\omega - \bar{\omega}_i)\left[\frac{\nu_{ei}}{\Omega_e}\frac{k_\perp^2}{k^2} - \frac{\Omega_e}{\nu_{ei}}\frac{k_z^2}{k^2} + i\frac{\hat{z} \times \underline{k}_\perp \cdot \nabla N_o}{N_o k^2}\right]$$

$$= -J_1^2(\xi)(\omega - \bar{\omega}_e)\frac{\Omega_i}{\omega_{pi}^2}\left[\frac{\omega_{pe}^2}{\omega - \bar{\omega}_i}\frac{1}{\Omega_i}\left(\frac{\nu_{in}}{\Omega_i}\frac{k_\perp^2}{k^2} + \frac{\Omega_i}{\nu_{in}}\frac{k_z^2}{k^2} + i\frac{\hat{z} \times \underline{k}_\perp \cdot \nabla N_o}{N_o k^2}\right)\right]$$

$$\times \left[1 + i\frac{\omega_{pe}^2}{\omega - \bar{\omega}_e}\frac{1}{\Omega_e}\left(\frac{\nu_{ei}}{\Omega_e}\frac{k_\perp^2}{k^2} - \frac{\Omega_e}{\nu_{ei}}\frac{k_\perp^2}{k^2} + i\frac{\hat{z} \times \underline{k}_\perp \cdot N_o}{N_o k^2}\right)\right]$$

$$\times \left[\frac{1}{H_e(\omega+\omega_o)} + \frac{1}{H_e(\omega-\omega_o)}\right] \tag{17}$$

Defining $\tilde{\omega}_{uH}^2 = \omega_{pe}^2 + \Omega_e^2 + \nu_e^2$, $\delta = \tilde{\omega}_{uH}^2 - \omega_o^2$, a measure of the frequency mismatch, $\delta_o = \frac{\nu_{ei}}{\omega_o}(\omega_{pe}^2 - 2\omega_o^2)$, we find from (17) (see Appendix 1) the modified growth rate

$$\gamma = \gamma_{\underline{k}}^L + 2\zeta^2\left(\frac{k_{\parallel}^2}{k_{\perp}^2} + \frac{\nu_{in}\nu_{ei}}{\Omega_i\Omega_e}\right)^{-1}\left(\frac{\hat{z}\times\underline{k}_{\perp}\cdot\nabla N_o}{k^2 N_o}\right)^2\left(\frac{\omega_{pe}}{\nu_{ei}}\frac{\nu_{ei}^2}{\Omega_e^2}\right)\frac{\omega_{pe}^2\delta}{\delta^2+\delta_o^2}\omega_{pe} \tag{18}$$

where
$$\zeta^2 = \frac{e^2}{m_e^2(\omega_o^2-\Omega_e^2)^2}\Big\{k_x^2 E_{px}^2 + k_y^2 E_{py}^2 + \frac{\Omega_e^2}{\omega_o^2}(k_y^2 E_{px}^2 + k_x^2 E_{py}^2)$$
$$+ 2E_{px}E_{py}\left[k_x k_y \cos\psi\left(1 - \frac{\Omega_e^2}{\omega_o^2}\right) - \frac{\Omega_e}{\omega_o}k_{\perp}^2\sin\psi\right]$$
$$+ \frac{\omega_o^2-\Omega_e^2}{\omega_o^2}k_z E_{pz}\left[k_x E_{px} - \frac{\Omega_e}{\omega_o}k_x E_{px}\sin\psi + k_y E_{py}\cos\psi\right]\Big\} \tag{19}$$

Here we have taken [Lee et al., 1972]

$$\underline{E}_p = \frac{E_{px}}{-2i}\exp(-i\omega_o t)\,\hat{\underline{x}} + \frac{E_{py}}{-2i}\exp(-i\omega_o t - i\psi)\hat{\underline{y}} + \frac{E_{pz}}{-2i}\exp(-i\omega_o t)\hat{\underline{z}} + \text{c.c.} \tag{20}$$

and included E_{pz} since we are considering modes for which $k_{\parallel} \neq 0$.

When $\delta > 0$, i.e., when $\omega_o < \tilde{\omega}_{uH}$ the pump induced term produces a destabilizing effect whereas when $\delta < 0$ it produces a stabilizing influence. We assume that the length of the irregularities along the magnetic field ($k_{\parallel} \ll k_{\perp}$) exceeds the length of regions in which destabilization ($\delta > 0$) or stabilization ($\delta < 0$) occurs. The pump induced term in eq. (18) contains terms proportional to ω_{pe}, an increasing (decreasing) function of altitude on the bottomside (topside) ionosphere, e.g., $\omega_{pe}^2(z) \simeq \omega_o^2(1 + z/L_N)$ with L_N constant, in local bottomside regions. In addition, Das and Fejer [1979] have shown that the altitude

dependence of the pump wave electric field near the resonance region must also be considered. As a result, a net destabilization (stabilization) will occur in bottomside regions if the pump electric field increases (decreases) with height in the resonance region. For topside regions, the opposite is true. For simplicity, in this discussion, we have neglected the effects of spatial dispersion. Spatial dispersion will cause a resonance at a somewhat lower height than when $\omega_o = (\omega_{pe}^2 + \Omega_e^2)^{1/2}$. The physical mechanism responsible for stabilization or destabilization of these convective instabilities can be understood as follows. A nonlinear interaction between $\underline{E} \times \underline{B}$ and current convective perturbation electric fields $\underline{\delta E}$ and the pump field $\underline{E}_p$ gives a new field $\underline{E}'$. The combined beating of the fields $\underline{E}'$ and $\underline{E}_p$ can produce a low frpequency force with component of the form $\exp[i(\underline{k}_L \cdot \underline{x} - \omega_L t + \delta)]$, $\underline{k}_L = \underline{k}' \pm \underline{k}_p$, $\omega_L = \omega' \pm \omega_p$, with phase determined by the sign of δ, which can in turn affect $\underline{\delta E}$ and modify the dispersion relation. This force will be pondermotive if $k_{L\parallel} > 1.55/\lambda_c$ whereas the partial pressure force will dominate if $k_{L\parallel} < 1.55/\lambda_c$ [Fejer, 1979]. Here λ_c is the electron mean free path and $k_{L\parallel}$ the component of $\underline{k}_L$ parallel to the magnetic field. We take $\underline{k} = k_x\hat{\underline{x}} + k_z\hat{z}$, an ambient electric field $E_o\hat{\underline{x}}$, and $N = N_o(y)$, i.e., a configuration unstable to the $\underline{E} \times \underline{B}$ and current convective instabilities. The growth rate in (18) is maximized for $\delta_{max} = \pm\,\delta_o = \pm\,(\nu_{ei}/\omega_o)\quad(\omega_{pe}^2 - 2\omega_o^2)$. Inserting this value for δ, eq. (18) can be written as

$$\gamma = \left(\frac{k_z^2}{k_x^2} + \frac{\nu_{in}}{\Omega_i}\frac{\nu_{ei}}{\Omega_e}\right)^{-1}\frac{\nu_{ei}}{\Omega_e}(k_xL)^{-1}\Big[\frac{\nu_{in}}{\Omega_i}\frac{c}{B}k_xE_o - k_zV_d \pm \frac{\zeta^2}{4}(k_xL)^{-1}$$

$$\times\left(\frac{\omega_{pe}}{\Omega_e}\right)\left(\frac{\omega_{pe}^2}{\nu_{ei}\omega_o}\right)\omega_{pe}\Big] - D_\perp k_\perp^2 - D_\parallel k_\parallel^2 \tag{21}$$

Here the upper (plus) sign refers to the case when $\delta > 0$ (destabilization) while the lower (minus) sign applies when $\delta < 0$ (stabilization). Eq. (21) gives the approximate minimum required power density for destabilization or stabilization

$$\left|\frac{cE_p^2}{8\pi}\right| = \frac{(k_xL)^2}{\pi}\frac{\omega_o}{\omega_{pe}^4}\left[(k_xL)^{-1}\frac{\nu_{ei}}{\Omega_e}\left(\frac{\nu_{in}}{\Omega_i}\frac{ck_xE_o}{B} - k_zV_d\right) - \left(D_\perp k_x^2 + D_\parallel k_z^2\right)\right.$$

$$\times \left(\frac{k_z^2}{k_x^2} + \frac{\nu_{in}}{\Omega_i}\frac{\nu_{ei}}{\Omega_e}\right)\Big] \frac{B^2(\omega_o^2 - \Omega_e^2)^2}{ck^2} \tag{22}$$

where we have assumed a circularly polarized O mode with $\psi = -\pi/2$ in eq. (19).

For typical ionospheric F region parameters at diffuse auroral latitudes in the 300-400 km altitude range, which give linear unstable growth ($\Omega_e \approx 8.5 \times 10^6 \sec^{-1}$, $\omega_{pe} \approx 3 \times 10^7 \sec^{-1}$, $\nu_{ei} \approx 5 \times 10^2 \sec^{-1}$, $\nu_{in}/\Omega_i \approx 10^{-4}$, $E_o \approx 10$ mV/m, $T_e \approx T_i \approx 1000^\circ$K, $j_\parallel \approx n_o e V_d \approx 1\ \mu A/m^2$, $kL \approx 2 \times 10^3$, $k^{-1} \approx 30$ m, $k_z/k_\perp \approx 10^{-4}$), we find from (22) the approximate power density for stabilization to be $cE_p^2/8\pi \approx 6 \times 10^{-4} W/m^2$. This is an overestimate since we have not considered electric field enhancements near the reflection point [Ginzburg, 1964]. For an linearly damped mode $k^{-1} \approx 5$m, we find, using eq. (22) the approximate power density necessary for destabilization ($\delta > 0$) to be $cE_p^2/8\pi = -3.2 \times 10^{-4} W/m^2$. This power density corresponds to a pump electric field amplitude of $E_p \approx 1.5 \times 10^{-2}$V/m. This electric field amplitude is comparable to that required by the thermal coupling mechanism of Vaskov and Gurevich [1975, 1977] and Das and Fejer [1979] to destabilize field-aligned density irregularities with scale lengths on the order of 0.5 - 5m. In reality, both the mechanism discussed in this paper and thermal mechanisms are probably acting simultaneously. From the coupling parameter ζ as given by (19) we see that there must be a component of $\underline{E}_p$ perpendicular to $\underline{B}$ since $k_\parallel/k_\perp \ll 1$. At high latitudes this can be accomplished using a vertically incident pump wave. An O-mode pump wave is preferred since the X-mode will be reflected at the right hand cutoff before it ever reaches the coupling region near the local upper hybrid frequency.

3. SUMMARY

The parametric effects of a large amplitude electromagnetic pump wave on convective plasma fluid instabilities ($\underline{E} \times \underline{B}$ gradient drift and current convective) in the high latitude F region ionosphere have been studied analytically. These convective instabilities have been invoked [Keskinen and Ossakow, 1982] to explain naturally occurring density irregularities in

and near large scale convecting auroral plasma enhancements [Vickrey et al., 1980]. We find that parametric coupling effects associated with a pump wave oscillating with approximately the upper hybrid frequency can stabilize or destabilize the $\underline{E} \times \underline{B}$ gradient drift and/or the current-convective instability. For parameters typical of the nighttime high latitude F region ionosphere, we find that these modes can be stabilized/destabilized with a free space incident power density on the order of 10^{-4}W/m^2. Since the pump wave should have a component of its electric field vector perpendicular to the magnetic field, a vertically propagating O mode is suggested. In addition, we note that these power density levels needed could be achieved using current ionospheric heaters located in Norway [Stubbe et al., 1981] and Alaska [Wong et al., 1981].

In our development we have made several approximations. We have assumed that the pump electric fields are spatially uniform (dipole approximation). We have ignored pump induced temperature and density changes. The time scale for these changes, to a first approximation, is longer than the parametric time scales of the instabilities studied here. In addition, we have neglected lower altitude absorption effects. Finally, we have assumed that the instabilities studied here are at or near marginal stability. We reserve these topics for a future report.

Appendix I

In solving eq. (17) for the growth rate V we must first evaluate the quantity $H_e^{-1}(\omega + \omega_o) + H_e^{-1}(\omega - \omega_o)$ where $H_e(\omega) = -[4\pi\, x_e(\omega)]^{-1}$ $[1 + 4\pi\, x_e(\omega)]$ and $x_e(\omega)$ is given by eq. (13). We have

$$H_e(\omega) = -\{i \frac{\omega_{pe}^2}{\omega-\bar{\omega}_e} [-\frac{1}{\check{\nu}_{ei}} \frac{k_z^2}{k^2} + \frac{1}{\Omega_e} \frac{\check{\kappa}_e}{1+\check{\kappa}_e} \frac{k_\perp^2}{k^2} + \frac{i}{\Omega_e} \frac{\check{\kappa}_e^2}{1+\check{\kappa}_e^2} \frac{\hat{z}\times \underline{k}_\perp \cdot \nabla N_o}{N_o k^2}]\}^{-1}$$

$$\times \{1 + i \frac{\omega_{pe}^2}{\omega-\bar{\omega}_e} [-\frac{1}{\check{V}_{ei}} \frac{k_z^2}{k^2} + \frac{1}{k^2} \frac{\check{\kappa}_e}{1+\check{\kappa}_e^2} \frac{k_\perp^2}{k^2} + \frac{i}{\Omega_e} \frac{\check{\kappa}_e^2}{1+\check{\kappa}_e} \frac{\hat{z}\times \underline{k}_\perp \cdot \nabla N_o}{N_o k^2}]\}$$

$$\simeq -[1 - i \frac{(\omega-\bar{\omega}_e)\Omega_e}{\omega_{pe}^2} \frac{1+\bar{\kappa}_e^2}{\check{\kappa}_e}]$$

$$= -[1 + \frac{\omega-\bar{\omega}_e}{\omega+i\nu_{ei}} \frac{\Omega_e^2-(\omega+i\nu_{ei})^2}{\omega_{pe}^2}] \qquad \text{(A1)}$$

where we have neglected the small k_z and ∇N_o effects on the sidebands $H_e(\omega+\omega_o)$ and $H_e(\omega-\omega_o)$. As a result we can write, with $\bar{\omega}_+ = \omega+\omega_o-\bar{\omega}_e$,

$$H_e^{-1}(\omega+\omega_o) = [1 + \frac{\bar{\omega}_+}{\bar{\omega}_++\bar{\omega}_e+i\nu_{ei}} \frac{\Omega_e^2-(\bar{\omega}_++\bar{\omega}_e+i\nu_{ei})^2}{\omega_{pe}^2}]^{-1}$$

$$= \frac{\bar{\omega}_++\bar{\omega}_e+i\nu_{ei}}{\omega_{pe}^2[\bar{\omega}_++\omega_e+i\nu_{ei}]+\bar{\omega}_+[\Omega_e^2-(\bar{\omega}_++\omega_e+i\nu_{ei})^2]}$$

$$= \frac{1+i\nu_{ei}(\bar{\omega}_++\bar{\omega}_e)^{-1}}{\omega_{pe}^2[1+i\nu_{ei}(\bar{\omega}_++\bar{\omega}_e)^{-1}]+\bar{\omega}_+(\bar{\omega}_++\bar{\omega}_e)^{-1}[\Omega_e^2-(\bar{\omega}_++\omega_e+i\nu_{ei})^2]}$$

$$\simeq \frac{1+i\nu_{ei}/\omega_o}{\delta + i\,\frac{\nu_{ei}}{\omega_o}\,(\omega_{pe}^2-2\omega_o^2}} \qquad \text{(A2)}$$

where we have assumed that $\bar{\omega}_+ + \omega_e \simeq \omega_o$, $\bar{\omega}_+ \simeq \omega_o$, $\nu_{ei}/\omega_o << 1$, and $\delta = \omega_{pe}^2 + \Omega_e^2 + \nu_e^2 - \omega_o^2$. Similarly,

$$H_e^{-1}(\omega-\omega_o) = \frac{1-i\nu_{ei}/\omega_o}{\delta - i\dfrac{\nu_{ei}}{\omega_o}(\omega_{pe}^2-2\omega_o^2)} \tag{A3}$$

As a result

$$H_e^{-1}(\omega+\omega_o) + H_e^{-1}(\omega-\omega_o) \simeq \delta(\delta^2+\delta_o^2)^{-1} \tag{A4}$$

where $\delta_o = (\nu_{ei}/\omega_o)(\omega_{pe}^2-2\omega_o^2)$. Inserting Eq. (A4) into eq. (17) it is straightforward to solve for $I_m\omega \equiv \gamma$ as given by eq. (18).

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